

Nuclear kaon dynamics

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Abstract

An effective low energy Lagrangian density is applied to nuclear K^- -dynamics. The free parameters, local s-wave couplings and $SU(3)$ -symmetry constrained range terms are adjusted to describe elastic and inelastic K^- -nucleon scattering data. The propagation and decay of the $\Lambda(1405)$ -resonance and the $\Lambda(1405)$ -nucleon hole state is studied self consistently with respect to the K^- -propagation in isospin symmetric nuclear matter.

1 Introduction

In this letter we consider kaon propagation in isospin symmetric nuclear matter. There has recently been much effort to evaluate the in medium K -mass in realistic models. The chiral Lagrangian has been applied perturbatively to the effective kaon mass [1,2]. Brown and Rho pursue a mean field approach in [3]. We endeavor a microscopic description deriving the kaon propagator in dense matter from kaon nucleon scattering in free space [4–8]. Consider the change of the K^+ -mass. As was emphasized in [10] it is given by the low density theorem in terms of the empirical K^+ -nucleon scattering lengths $a_{K+N}^{(0)} \simeq 0.02$ fm and $a_{K+N}^{(1)} \simeq -0.32$ fm [9]

$$\Delta m_K^2 = -\pi \left(1 + \frac{m_K}{m_N}\right) \left(a_{KN}^{(I=0)} + 3 a_{KN}^{(I=1)}\right) \rho + \mathcal{O}(k_F^4) \quad (1)$$

where $\rho = 2 k_F^3 / (3 \pi^2)$. Model calculations [2,7] typically find only small corrections to (1). In fact the next to leading term of order k_F^4 can be evaluated exclusively in terms of the K^+N -scattering lengths. We obtain the model independent result

$$\Delta m_K^2 = (1) + \alpha \left(\left(a_{K+N}^{(I=0)}\right)^2 + 3 \left(a_{K+N}^{(I=1)}\right)^2 \right) k_F^4 + \mathcal{O}(k_F^5) \quad (2)$$

where

$$\alpha = \frac{1 - x^2 + x^2 \log(x^2)}{\pi^2 (1 - x)^2} \simeq 0.166 \quad (3)$$

and $x = m_K/m_N$. The correction term is indeed small. At nuclear saturation density with $k_F \simeq 265$ MeV it increases the repulsive K^+ -mass shift from 28 MeV to 35 MeV by about 20%. Thus the density expansion is useful in the K^+ -channel. Any microscopic model consistent with low energy K^+ -nucleon scattering data is bound to give similar results for the K^+ -propagation in nuclear matter at densities $\rho \simeq \rho_0$ sufficiently small to maintain the density expansion rapidly convergent. Consequently the role of chiral symmetry is restricted to the qualitative prediction of the K^+N -scattering lengths by the Weinberg-Tomozawa term.

We continue with the K^- -mass. Again as a first step one may apply the low density theorem. The empirical scattering lengths $a_{K^-N}^{(0)} \simeq (-1.70 + i 0.68)$ fm and $a_{K^-N}^{(1)} \simeq (0.37 + i 0.60)$ fm [9,11] imply according to (1) a repulsive mass shift of 23 MeV with a width of $\Gamma_{K^-} \simeq 147$ MeV at saturation density. The correction term (2) results in a total repulsive mass shift of 55 MeV and a width of $\Gamma_{K^-} \simeq 195$ MeV. At nuclear saturation the density expansion for the K^- -mode is poorly convergent if at all. Furthermore, the leading terms appear to contradict kaonic atom data [12] which suggest sizable attraction at small density. Finally the empirical K^-N scattering lengths are in striking disagreement with the Weinberg-Tomozawa term, the leading order chiral prediction.

The solution to this puzzle lies in the presence of the $\Lambda(1405)$ resonance in the K^- -proton channel [5,6,13]. The $\Lambda(1405)$ -resonance can be described together with elastic and inelastic K^- -proton scattering data in terms of a coupled channel Lippman-Schwinger equation with the potential matrix evaluated perturbatively from the chiral Lagrangian [13]. We also note that a satisfactory description of kaon nucleon scattering data can be achieved by the coupled K -matrix approach of Martin [9].

Contrary to the K^+ -propagation in nuclear matter different microscopic models consistent with low energy kaon nucleon scattering data may predict different results for the K^- -propagation in nuclear matter simply because an attractive in-medium K^- -self energy probes the K^- - nucleon scattering amplitude below the kaon nucleon threshold. Below the physical threshold the amplitude is subject to uncertainties due to the necessary subthreshold extrapolation of scattering data. Also, as was pointed out first by Koch [5], the $\Lambda(1405)$ resonance may experience a repulsive mass shift due to Pauli blocking which strongly affects the in-medium K^- -nucleon scattering amplitude [5,6]. This offers a simple mechanism for the transition from repulsion, implied by

the scattering lengths, at low densities, $\rho < 0.1 \rho_0$, to attraction at somewhat larger densities [6] as favored by kaonic atom data. Schematically this effect can be reproduced in terms of an elementary $\Lambda(1405)$ field dressed by a kaon nucleon loop. The repulsive Λ mass shift due to the Pauli blocking of the nucleon is given by:

$$\Delta m_\Lambda = \frac{g_{\Lambda NK}^2}{\pi^2} \frac{m_N}{m_\Lambda} \left(1 - \frac{\mu_\Lambda}{k_F} \arctan \left(\frac{k_F}{\mu_\Lambda} \right) \right) k_F \quad (4)$$

with the 'small' scale

$$\mu_\Lambda^2 = \frac{m_N}{m_\Lambda} \left(m_K^2 - (m_\Lambda - m_N)^2 \right) \simeq (144 \text{ MeV})^2 \quad (5)$$

and the $\Lambda(1405)$ kaon nucleon coupling constant $g_{\Lambda NK}$.

In this letter we extend previous work [5,7] and treat the K^- -state and the $\Lambda(1405)$ states self consistently. The $\Lambda(1405)$ -resonance mass in matter is then the result of two competing effects: the Pauli blocking increases the mass whereas the decrease of the K^- mass tends to lower the mass since the $\Lambda(1405)$ can be considered as a K^- -proton bound state. We expect this mechanism to be important for in-medium K^-N -scattering simply because the characteristic scale μ_Λ in (5) depends sensitively on small variations of m_Λ and m_K .

2 Kaon nucleon scattering

We describe K^- nucleon scattering by means of an effective Lagrangian density. Consider first the isospin zero ($I=0$) channel:

$$\begin{aligned} \mathcal{L} = & \frac{1}{2} g_{11}^{(I=0)} (N^\dagger K) (K^\dagger N) + \frac{1}{\sqrt{6}} g_{12}^{(I=0)} (N^\dagger K) (\vec{\pi}^\dagger \cdot \vec{\Sigma}) \\ & + \frac{1}{\sqrt{6}} g_{21}^{(I=0)} (\vec{\Sigma}^\dagger \cdot \vec{\pi}) (K^\dagger N) + \frac{1}{3} g_{22}^{(I=0)} (\vec{\Sigma}^\dagger \cdot \vec{\pi}) (\vec{\pi}^\dagger \cdot \vec{\Sigma}) \end{aligned} \quad (6)$$

with the isospin doublet fields $K = (K_-^\dagger, \bar{K}_0^\dagger)$ and $N = (p, n)$. Here we include the pion and the $\Sigma(1195)$ as relevant degrees of freedom since they couple strongly to the K^- nucleon system. The nucleon and Σ as well as the kaon and pion fields are constructed with relativistic kinematics but without anti-particle components. The free nucleon and kaon propagators take the form:

$$S_N(\omega, \vec{q}) = \frac{m_N}{E_N(q)} \frac{1}{\omega - E_N(q) + i\epsilon}$$

$$S_K(\omega, \vec{q}) = \frac{1}{2 E_K(q)} \frac{1}{\omega - E_K(q) + i\epsilon}, \quad (7)$$

respectively, where $E_a(q) = \sqrt{m_a^2 + q^2}$. The isospin zero coupled channel scattering amplitude

$$T = \begin{pmatrix} T_{KN \rightarrow KN} & T_{KN \rightarrow \pi\Sigma} \\ T_{\pi\Sigma \rightarrow KN} & T_{\pi\Sigma \rightarrow \pi\Sigma} \end{pmatrix} \quad (8)$$

is given by the set of ladder diagrams resummed conveniently in terms of the Bethe-Salpeter integral equation. Since the interaction terms in (6) are local the Bethe-Salpeter equation reduces to the simple matrix equation

$$T(s) = g(s) + g(s) J(s) T(s) = \left(g^{-1}(s) - J(s) \right)^{-1}. \quad (9)$$

with the loop matrix $J = \text{diag} (J_{KN}, J_{\pi\Sigma})$ and

$$\begin{aligned} J_{KN}(\omega, \vec{q}) &= - \int_0^\lambda \frac{d^3 l}{(2\pi)^3} \frac{m_N}{E_N(l)} S_K(\omega - E_N(l), \vec{q} - \vec{l}) \\ J_{\pi\Sigma}(\omega, \vec{q}) &= - \int_0^\lambda \frac{d^3 l}{(2\pi)^3} \frac{m_\Sigma}{E_\Sigma(l)} S_\pi(\omega - E_\Sigma(l), \vec{q} - \vec{l}). \end{aligned} \quad (10)$$

Small range terms are included in our scheme by the replacements $g_{11} \rightarrow g_{11} + h_{11} (s - (m_N + m_K)^2)$, $g_{12} \rightarrow g_{12} + h_{12} (s - (m_N + m_K)^2)$ and $g_{22} \rightarrow g_{22} + h_{22} (s - (m_\Sigma + m_\pi)^2)$ induced by appropriate additional terms in (6). The loop functions J_{KN} and $J_{\pi\Sigma}$ are regularized by the cutoff $\lambda = 0.7$ GeV. For small three momenta $|\vec{q}| < \lambda$ the loop functions $J_{KN}(\omega, \vec{q})$ and $J_{\pi\Sigma}(\omega, \vec{q})$ depend to good accuracy exclusively on the combination $s = \omega^2 - \vec{q}^2$ as expected from covariance.

The Lagrangian density (6) follows from a chiral Lagrangian with relativistic baryon and meson fields upon integrating out the anti-particle field components. Therefore the coupling matrix g is constrained to some extent by chiral symmetry [13]. To leading order we derive the isospin zero coupling strengths

$$g^{(I=0)} = \begin{pmatrix} 3 \frac{m_K}{2 f_\pi^2} & \sqrt{6} \frac{m_\pi + m_K}{8 f_\pi^2} \\ \sqrt{6} \frac{m_\pi + m_K}{8 f_\pi^2} & 2 \frac{m_\pi}{f_\pi^2} \end{pmatrix} \quad (11)$$

by matching tree level threshold amplitudes. The chiral matching of correction terms and the range parameters h_{ij} is less obvious and not pursued here.

In fact a consistent chiral matching requires the K^- -nucleon potential to be evaluated minimally at chiral order Q^3 . Only at this order the required counter terms for the loop functions are introduced. In this work the coupling strengths g_{ij} and the range parameters h_{ij} are directly adjusted to reproduce empirical scattering data described in terms of the coupled channel scattering amplitude (9). The set of parameters $g_{11} \lambda = 46.86$, $g_{12} \lambda = 11.67$, $g_{22} \lambda = 16.08$, $h_{11} \lambda^3 = 0.79$, $h_{12} \lambda^3 = 8.57$ and $h_{22} \lambda^3 = 4.94$ results from a least square fit to the amplitudes of [13]. We obtain a good description of all coupled channel amplitudes with the isospin zero scattering length $a_{K^-N}^{(I=0)} \simeq (-1.76 + i 0.60)$ fm. Our parameters confirm the result of [13] that the Weinberg-Tomozawa term (11) predicts the interaction strength in the various channels rather accurately. Fig. 1 shows that the isospin zero scattering amplitude $f_{KN}^{(I=0)}(\omega) = m_N T_{KN}^{(I=0)}(\omega, \vec{q} = 0)/(4\pi\omega)$ is clearly dominated by the $\Lambda(1405)$ resonance.

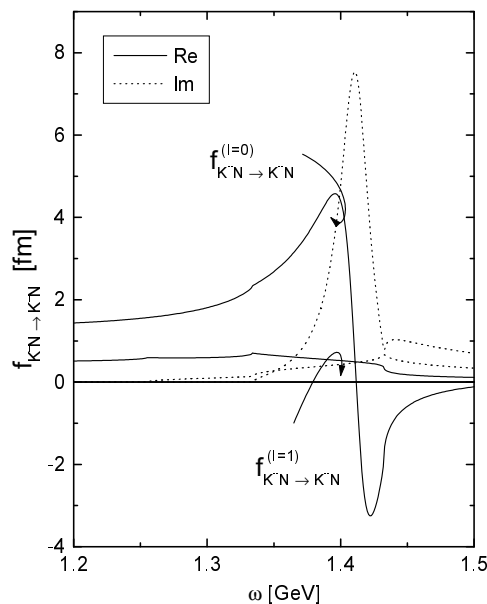


Fig. 1. K^- -nucleon scattering amplitude.

Let us now turn to the $I = 1$ channel. Here the K^- nucleon system couples strongly also to the $\pi \Lambda(1115)$ -channel. The appropriate effective Lagrangian density is:

$$\begin{aligned} \mathcal{L} = & \frac{1}{2} g_{11}^{(I=1)} (N^\dagger \vec{\tau} K) (K^\dagger \vec{\tau} N) - \frac{1}{2} g_{22}^{(I=1)} (\vec{\Sigma}^\dagger \times \vec{\pi}) (\vec{\pi}^\dagger \times \vec{\Sigma}) \\ & + g_{33}^{(I=1)} (\Lambda^\dagger \vec{\pi}) (\vec{\pi}^\dagger \Lambda) - \frac{i}{2} g_{12}^{(I=1)} [(N^\dagger \vec{\tau} K) (\vec{\pi}^\dagger \times \vec{\Sigma}) - h.c.] \\ & + \frac{1}{\sqrt{2}} g_{13}^{(I=1)} [(N^\dagger \vec{\tau} K) (\vec{\pi}^\dagger \Lambda) + h.c.] \end{aligned}$$

$$-\frac{i}{\sqrt{2}} g_{23}^{(I=1)} \left[\left(\vec{\Sigma}^\dagger \times \vec{\pi} \right) \left(\vec{\pi}^\dagger \Lambda \right) - h.c. \right] \quad (12)$$

We construct the $I = 1$ coupled channel scattering amplitude in full analogy to the isospin zero case with $J = \text{diag} (J_{KN}, J_{\pi\Sigma}, J_{\pi\Lambda})$ (see eq. (9)). The range terms are included by the replacements $g_{11} \rightarrow g_{11} + h_{11} (s - (m_N + m_K)^2)$, $g_{12} \rightarrow g_{12} + h_{12} (s - (m_N + m_K)^2)$, $g_{13} \rightarrow g_{13} + h_{13} (s - (m_N + m_K)^2)$, $g_{22} \rightarrow g_{22} + h_{22} (s - (m_\Sigma + m_\pi)^2)$ and $g_{33} \rightarrow g_{33} + h_{33} (s - (m_\Lambda + m_\pi)^2)$. The coupling strengths as obtained by chiral matching of tree level threshold amplitudes are:

$$g^{(I=1)} = \begin{pmatrix} \frac{m_K}{2 f_\pi^2} & \frac{m_\pi + m_K}{4 f_\pi^2} \sqrt{6} \frac{m_\pi + m_K}{8 f_\pi^2} \\ \frac{m_\pi + m_K}{4 f_\pi^2} & \frac{m_\pi}{f_\pi^2} & 0 \\ \sqrt{6} \frac{m_\pi + m_K}{8 f_\pi^2} & 0 & 0 \end{pmatrix}. \quad (13)$$

We point out that the range terms can be expressed in terms of the isospin zero range parameters and one free parameter h_F

$$\begin{aligned} h_{11}^{(I=1)} &= \frac{1}{2} \left(h_{11}^{(I=0)} - \sqrt{6} h_{12}^{(I=0)} + h_{22}^{(I=0)} - 6 h_F \right) \\ h_{12}^{(I=1)} &= \frac{1}{6} \left(3 h_{11}^{(I=0)} + \sqrt{6} h_{12}^{(I=0)} - 3 h_{22}^{(I=0)} - 6 h_F \right) \\ h_{13}^{(I=1)} &= \frac{1}{12\sqrt{6}} \left(6 h_{11}^{(I=0)} + 2 \sqrt{6} h_{12}^{(I=0)} - 6 h_{22}^{(I=0)} + 36 h_F \right) \\ h_{22}^{(I=1)} &= \frac{1}{12} \left(15 h_{11}^{(I=0)} - 5 \sqrt{6} h_{12}^{(I=0)} - 3 h_{22}^{(I=0)} - 30 h_F \right) \\ h_{23}^{(I=1)} &= 0 \\ h_{33}^{(I=1)} &= \frac{1}{9} \left(3 h_{11}^{(I=0)} - 5 \sqrt{6} h_{12}^{(I=0)} + 6 h_{22}^{(I=0)} - 18 h_F \right) \end{aligned} \quad (14)$$

if constrained by $SU(3)$ -symmetry. We obtain a good description of the isospin one amplitudes $K^- N \rightarrow K^- N, \pi\Sigma, \pi\Lambda$ of [13] with $g_{23} = g_{33} = 0$ as suggested by the leading order result (13). The K^- -nucleon scattering amplitude $f_{K^-N}^{(I=1)}(\omega) = m_N T_{K^-N}^{(I=1)}(\omega, \vec{q} = 0)/(4\pi\omega)$ as shown in Fig. 1 follows with the set of parameters $g_{11}\lambda = 12.75$, $g_{12}\lambda = 13.56$, $g_{13}\lambda = 15.07$, $g_{22}\lambda = 16.04$ and $h_F\lambda^3 = -1.92$. The scattering length comes at $a_{K^-N}^{(I=1)} \simeq (0.35 + i 0.69)$ fm. We observe that the Weinberg-Tomozawa term (13) predicts the interaction strengths in the $I = 1$ channel less accurately than in the $I = 0$ channel. Furthermore, the imposed $SU(3)$ -symmetry for the range parameters (14) is found to be essential for our subthreshold extrapolation of the $K^- N$ -scattering amplitude in the $I = 1$ channel.

3 Kaon self energy in nuclear matter

The kaon self energy $\Pi(\omega, \vec{q})$ is evaluated in the nucleon gas approximation. Its imaginary part follows from the imaginary part of the in medium kaon nucleon scattering amplitude \bar{T}_{K-N} :

$$\Im \Pi_K(\omega, \vec{q}) = -4 \int_0^{k_F} \frac{d^3 l}{(2\pi)^3} \frac{m_N}{E_N(l)} \Im \bar{T}_{KN}(\omega + E_N(l), \vec{l} + \vec{q}) \cdot \Theta \left(l^2 - k_F^2 - \omega^2 + 2\omega \sqrt{m_N^2 + k_F^2} \right) \quad (15)$$

with the Fermi momentum k_F . The θ -function emerges from the zero temperature limit of appropriate Fermi-Dirac distribution functions (see e.g. [14]) and ensures that the kaon spectral density vanishes at zero energy. The real part of the self energy then is evaluated by means of the dispersion integral:

$$\Pi_K(\omega, \vec{q}) = \int_{-\infty}^{\infty} \frac{d\bar{\omega}}{\pi} \frac{\Im \Pi_K(\bar{\omega}, \vec{q})}{\bar{\omega} - \omega - i\epsilon} \quad (16)$$

The kaon nucleon scattering amplitude $4\bar{T}_{KN} = \bar{T}_{KN}^{(I=0)} + 3\bar{T}_{KN}^{(I=1)}$ is medium modified exclusively through the kaon nucleon loop since it is given by (9) with the vacuum kaon nucleon loop J_{KN} replaced by the in matter loop \bar{J}_{KN} . Selfconsistency is met once \bar{J}_{KN} is evaluated in terms of the kaon propagator:

$$\bar{S}_K(\omega, \vec{q}) = \frac{1}{2 E_K(q)} \frac{1}{\omega - E_K(q) - \Pi(\omega, \vec{q})/(2 E_K(q)) + i\epsilon} \quad (17)$$

with the kaon self energy $\Pi_K(\omega, \vec{q})$ of (15). In the nucleon gas approximation the loop function reads:

$$\Im \bar{J}_{KN}(\omega, \vec{q}) = - \int_{k_F}^{\infty} \frac{d^3 l}{(2\pi)^3} \frac{m_N}{E_N(l)} \Im \bar{S}_K(\omega - E_N(l), \vec{q} - \vec{l}) \cdot \Theta \left(l^2 - \omega^2 - m_N^2 \right) \quad (18)$$

with the real part determined by the dispersion integral:

$$\bar{J}_{KN}(\omega, \vec{q}) = \int_{-\infty}^{\infty} \frac{d\bar{\omega}}{\pi} \frac{\Im \bar{J}_{KN}(\bar{\omega}, \vec{q})}{\bar{\omega} - \omega - i\epsilon} \quad (19)$$

It is convenient to formally rewrite the set of coupled equations as follows:

$$\begin{aligned}\bar{J}_{KN}(\omega, \vec{q}) &= - \int_{k_F}^{\lambda} \frac{d^3 l}{(2\pi)^3} \frac{m_N}{E_N(l)} \bar{S}_K(\omega - E_N(l), \vec{q} - \vec{l}) + \Delta \bar{J}_{KN}(\omega, \vec{q}) \\ \Delta \bar{J}_{KN}(\omega, \vec{q}) &= \int_{-\infty}^{m_N} \frac{d\bar{\omega}}{\pi} \int_{\sqrt{\bar{\omega}^2 - m_N^2}}^{\infty} \frac{d^3 l}{(2\pi)^3} \frac{m_N}{E_N(l)} \frac{\Im \bar{S}_K(\bar{\omega} - E_N(l), \vec{q} - \vec{l})}{\bar{\omega} - \omega - i\epsilon}\end{aligned}\quad (20)$$

with the kaon nucleon loop, $\bar{J}_{KN}(\omega, \vec{q})$, now regularized by our cutoff parameter λ such as to reproduce the vacuum loop function $J_{KN}(\omega, \vec{q})$ in the zero density limit. Similarly we write

$$\begin{aligned}\Pi_K(\omega, \vec{q}) &= -4 \int_0^{k_F} \frac{d^3 l}{(2\pi)^3} \frac{m_N}{E_N(l)} \bar{T}_{KN}(\omega + E_N(l), \vec{q} + \vec{l}) + \Delta \Pi_K(\omega, \vec{q}) \\ \Delta \Pi_K(\omega, \vec{q}) &= 4 \int_{-\infty}^{\mu - m_N} \frac{d\bar{\omega}}{\pi} \int_0^{\sqrt{k_F^2 - 2\bar{\omega}\mu}} \frac{d^3 l}{(2\pi)^3} \frac{m_N}{E_N(l)} \frac{\Im \bar{T}_{KN}(\bar{\omega} + E_N(l), \vec{q} + \vec{l})}{\bar{\omega} - \omega - i\epsilon}\end{aligned}\quad (21)$$

with the nucleon chemical potential $\mu^2 = m_N^2 + k_F^2$. It is immediate that $\Im \Delta \Pi_K(\omega, \vec{q}) = 0$ for $\omega > \sqrt{m_N^2 + k_F^2} - m_N$ and $\Im \Delta \bar{J}_{KN}(\omega, \vec{q}) = 0$ for $\omega > m_N$. Moreover suppose $\Im T_{KN}(\omega) = 0$ for $\omega < \omega_{thres.}$. The free space amplitude of our model suggests $\omega_{thres.} = m_\Lambda + m_\pi$. Then $\Delta \Pi_K = 0$ provided that $\mu < \omega_{thres.}$ or $k_F < \sqrt{(m_\Lambda + m_\pi)^2 - m_N^2} \simeq 830$ MeV holds. Similar one finds $\Delta \bar{J}_{KN} = 0$ since $\mu < \omega_{thres.}$ implies also $\Im \Pi_K(\omega) = 0$ for $\omega < 0$. Hence it is justified to solve our coupled set of integral equations with $\Delta \Pi_K = \Delta \bar{J}_{KN} = 0$.

We compare various degrees of approximations for the kaon self energy. The set of equations for the kaon self energy and the scattering amplitude is solved iteratively. We find that self consistency is reached after few iterations: the kaon nucleon loop function, \bar{J}_{KN} , is approximated surprisingly well if evaluated with the kaon spectral density as obtained from the free space scattering amplitude. Fig. 2 shows contour plots of the kaon spectral density as evaluated with the free space kaon nucleon scattering amplitude, with the Pauli blocked amplitude (the approximation scheme applied in [5–7]) and with the self consistent amplitude. All three schemes predict a two mode structure of the spectral density, however, with quantitative differences. Both, the Pauli blocked and the self consistent amplitude shift strength from the upper branch to the lower branch as compared with the spectral density derived from the free space amplitude. While the Pauli blocked amplitude causes a small repulsive

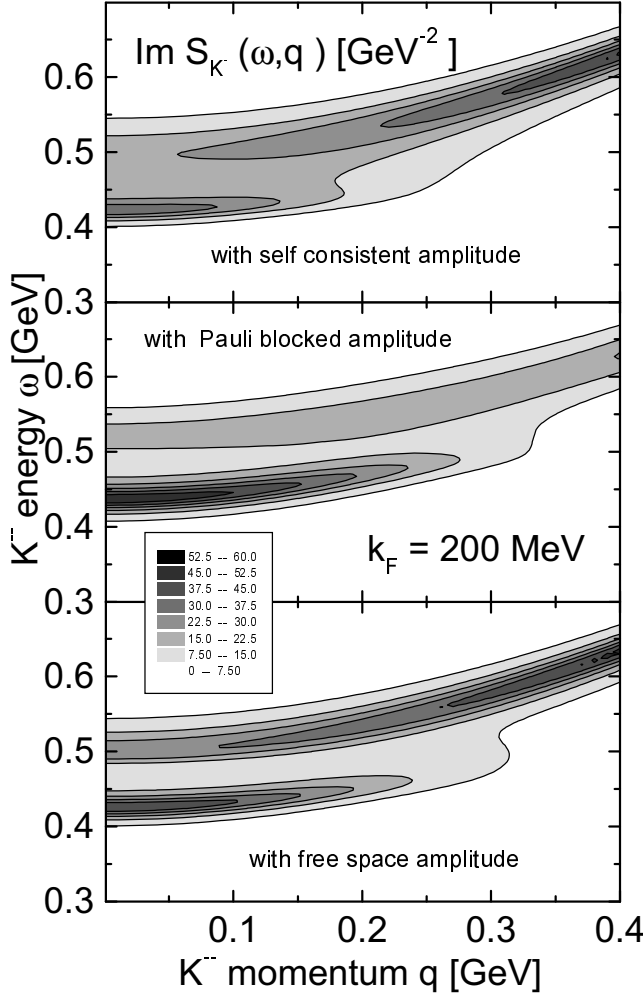


Fig. 2. Contour plot of K^- -spectral density at $k_F = 200$ MeV.

shift of the $\Lambda(1405)$ -nucleon hole state the selfconsistent amplitude predicts an attractive shift. At larger densities self consistency affects the K^- -rest mass moderately. For example at $k_F = 300$ MeV we find $\Delta m_{K^-} \simeq -140$ MeV and $\Gamma_{K^-} \simeq 35$ MeV as compared with $\Delta m_{K^-} \simeq -141$ MeV and $\Gamma_{K^-} \simeq 31$ MeV from the free space amplitude and $\Delta m_{K^-} \simeq -123$ MeV and $\Gamma_{K^-} \simeq 29$ MeV from the Pauli blocked amplitude. Note here that the quasi particle width $\Gamma_{K^-} = -\Im \Pi(m_{K^-} + \Delta m_{K^-}, \vec{q} = 0)/(m_{K^-} + \Delta m_{K^-})$, given above, differs from the physical K^- -width by about a factor of two due to the strong energy dependence of the kaon self energy (see Fig. 3). We find that self consistency is most important once the K^- -mode starts moving relative to the nuclear medium. Here we expect p-wave K^-N -interactions [15], which are not included in this work, to modify our results to some extent.

In Fig. 3 we present our final result for the kaon spectral density as a function of the kaon energy ω for various Fermi momenta k_F and kaon momenta \vec{q} . Typ-

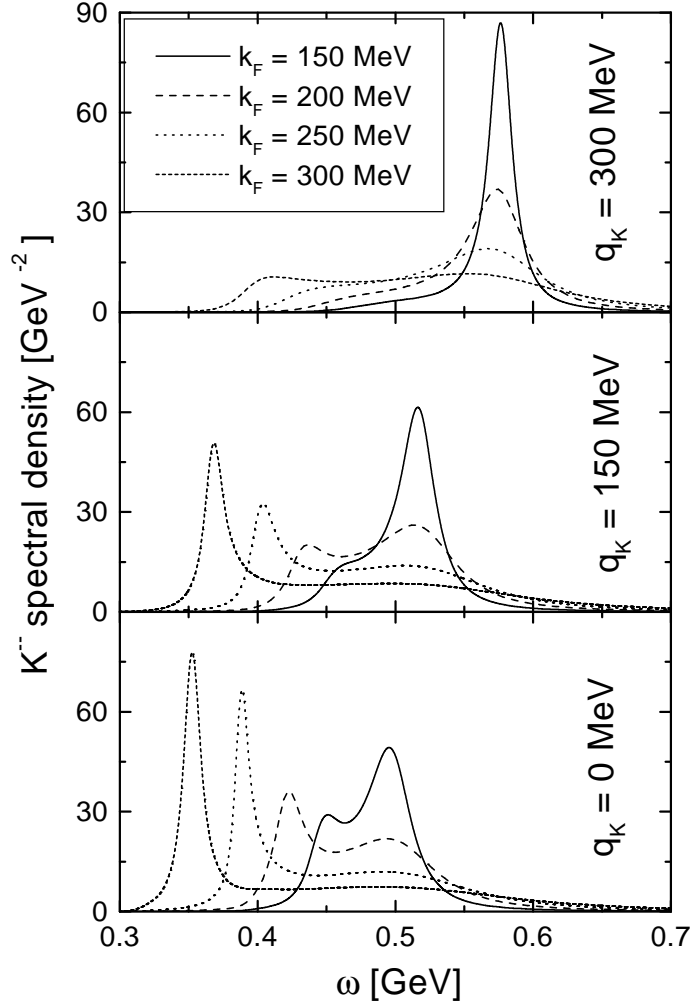


Fig. 3. K^- spectral density.

ically the spectral density exhibits a two peak structure representing the K^- and the $\Lambda(1405)$ -nucleon hole states. As the Fermi momentum k_F increases the energetically lower state experiences a strong attractive shift whereas the more massive state becomes broader. On the other hand as the kaon momentum increases both states basically gain kinetic energy with the energetically higher peak attaining more strength. We find that the sum rule for the kaon spectral density

$$\int_0^\infty d\omega \Im S_K(\omega, \vec{q}) = \frac{\pi}{2\sqrt{m_K^2 + \vec{q}^2}} \quad (22)$$

holds at the 5% accuracy level.

Fig. 4 shows the resulting isospin zero K^- -nucleon scattering amplitudes with

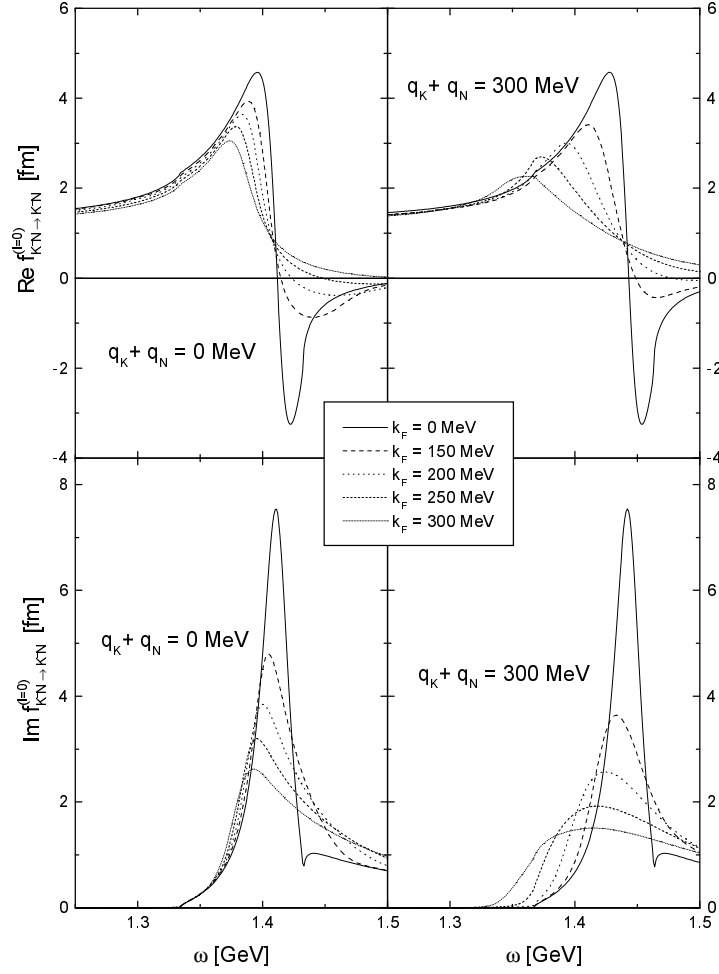


Fig. 4. K^- -nucleon scattering in nuclear matter.

$|\vec{q}_K + \vec{q}_N| = 0$ and $|\vec{q}_K + \vec{q}_N| = 300$ MeV. The imaginary part of the amplitude shows a clear peak around 1.4 GeV for all Fermi momenta representing the $\Lambda(1405)$ resonance state. As nuclear matter is compressed the peak gets broader with little effect on the peak position. The real part of the scattering amplitude changes its sign for $\omega > m_N + m_K$ from repulsion to attraction as the density gets larger.

As argued in the introduction the $\Lambda(1405)$ -mass shift results from the repulsive Pauli blocking effect and the attractive feedback effect of a decreased kaon mass. Altogether the $\Lambda(1405)$ resonance mass is left more or less at its free space value, however, with an increased decay width.

References

- [1] G.E. Brown, K. Kubodera, M. Rho, V. Thorsson, Phys. Lett. **B 291** 355, (1992);
G.E. Brown, C.H. Lee, M. Rho, V. Thorsson, Nucl. Phys. **A 567** 937, (1994);
C.-H. Lee, G.E. Brown, D.-P. Min, M. Rho, Nucl. Phys. **A 385**, 481 (1995).
- [2] C.-H. Lee, D.-P. Min and M. Rho, Nucl. Phys. **A 602**, 334 (1996).
- [3] G.E. Brown and M. Rho, Nucl. Phys. **A 596**, 503 (1996); G.Q. Li, C.-H. Lee,
G.E. Brown, nucl-th/9706057.
- [4] H. Yabu et al. Phys. Lett. **B 315**, 17 (1993).
- [5] V. Koch, Phys. Lett. **B 337**, 7 (1994); A. Ohnishi, Y. Nara and V. Koch, nucl-
th/9706084.
- [6] T. Waas, N. Kaiser and W. Weise, Phys. Lett. **B 365**, 12 (1996), Phys. Lett. **B**
379, 34 (1996).
- [7] T. Waas, M. Rho and W. Weise, Nucl. Phys. **A 617**, 449 (1997).
- [8] V.R. Pandharipande, C.J. Pethick and V. Thorsson, Phys. Rev. Lett. **75**, 4567
(1995).
- [9] A.D. Martin, Nucl. Phys. **B 179**, 33 (1981).
- [10] M. Lutz, A. Steiner and W. Weise, Nucl. Phys. **A 574**, 755 (1994).
- [11] M. Iwasaki et al., Phys. Rev. Lett. **78**, 3067 (1997).
- [12] E. Friedmann, A. Gal and C.J. Batty, Nucl. Phys. **A 579** , 518 (1994).
- [13] N. Kaiser, P.B. Siegel and W. Weise, Nucl. Phys. **A 594**, 325 (1995); N. Kaiser,
T. Waas and W. Weise, Nucl. Phys. **A 612**, 297 (1997); T. Waas, private
communication.
- [14] C.L. Korpa and R. Malfliet, Phys. Rev. **C 52**, 2756 (1995)
- [15] E.E. Kolomeitsev, D.N. Voskresensky, B. Kämpfer, Nucl. Phys **A 588** 889,
(1995).